

On Field-Induced Quantum Criticality in $YbRh_2Si_2$.

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The quantum critical point (QCP) in $YbRh_2Si_2$ is an enigma for the itinerant view of QCP. In an alternative view, this QCP is intimately linked to the selective Mott localization of the heavy f electrons. Following a perusal of this unusual QCP, I study an Extended Periodic Anderson Model (EPAM) within DMFT. A quantum phase transition (FQPT), accompanied by a rapid change in the Fermi volume, is found near the quantum-critical end-point of the selective Mott transition in the f -electron sector. The theory accounts for a wide range of unusual, singular non-Fermi liquid features exhibited at this QCP in $YbRh_2Si_2$ in a natural way.

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Quantum Criticality in solids and the potential to “tune in” to new, novel phases of matter in their proximity underpins a large component of modern condensed matter physics research [1, 2]. The rare-earth compound $YbRh_2Si_2$ is a particularly enigmatic case in point, exhibiting a “strange” non-FL “phase” fanning out from a $T = 0$ QCP separating antiferromagnetic (AF) ordered and heavy Fermi liquid (HFL) phases. $YbRh_2Si_2$ is tuned to the QCP by minute magnetic field or chemical substitution. Experimental data indicate that the Hertz-Moriya-Millis (HMM) scenario [3] does not account for the unusual responses in $YbRh_2Si_2$. An alternative view posits that these are associated with the destruction of the Kondo effect itself, either by inter-site RKKY interactions [2, 5], or by selective localization of f electrons [6]. Given that the “standard model” of f -band systems, the Periodic Anderson Model (PAM), emphasizes the HFL aspect driven by *quasilocal* Kondo screening [7], these new observations call for mechanisms which destabilize Kondo singlet formation. In spite of vigorous attempts [2, 6], the problem is far from “being solved”.

I begin by recapitulating salient features of the FQCP in $YbRh_2Si_2$:

(i) dc resistivity, $\rho(T) \simeq AT$ over three decades in T [8], (ii) specific heat, $C_V(T) \simeq T^{0.6}$ at very low T [2], (iii) anomalously slow (in frequency, ω) decay of optical conductivity, with linear-in- ω scattering rate, $\tau^{-1}(\omega) \simeq \omega^\alpha$ with $\alpha \simeq 1$ [9], (iv) strongly T -dependent Hall constant, $R_H(T)$, and $\cot\theta_H(T) \simeq C_1 T^2 + C_2$, and (v) rapid change in the low- T value of R_H across the FQCP, extrapolating to a jump as $T \rightarrow 0$ [10], suggesting a rapid change in the Fermi surface (FS) across the FQCP, (vi) static magnetic susceptibility, $\chi(\mathbf{q} = 0, T, B = 0) \simeq T^{-0.6}$ for $T > 0.3$ K, along with large Korringa ratio [12], indicating very strong *ferromagnetic* correlations close to the FQCP. And $\chi(B) \simeq (B - B_c)^{-0.6}$ scales with the A -coefficient of the T^2 term in $\rho(T)$ in the HFL regime, (vii) NMR derived Knight shift, $K_s(T, B)$ and the relaxation rate, $1/T_1 T$ scale with $\chi(\mathbf{q} = 0, T, B = 0)$.

(i),(iii) and (iv) are reminiscent of what is seen in high-

T_c cuprates in their “normal” state. Recently, (v) has also been seen in cuprates near optimal doping [13]. All these behaviors are at odds with the HMM theory [3], which predicts markedly different behavior [4]. So the elucidation of (i)-(vii) in a single theoretical picture remains a challenge.

Here, I address these issues by proposing a modified PAM with extended f -hopping and hybridisation, as well as a direct coulomb interaction between the f electrons and conduction (c) electrons, dubbed Extended-PAM (EPAM). I study this EPAM using DMFT, showing how the non-FL behavior along a curve in parameter space is understood as a selective Mott localization, and discuss how (i)-(vii) naturally follow therefrom. To the extent that this non-FL behavior is tied to f -Mott physics, single-site DMFT should capture the relevant physics. I will also show how this non-FL state is unstable to either AF, or to a heavy FL (HFL) away from this curve, at $T = 0$.

The Hamiltonian is $H = H_0 + H_1$, with the band part described by

$$H_0 = -t_f \sum_{\langle i,j \rangle, \sigma} f_{i\sigma}^\dagger f_{j\sigma} - t_p \sum_{\langle i,j \rangle, \sigma} c_{i\sigma}^\dagger c_{j\sigma} + V_{fc} \sum_{\langle i,j \rangle, \sigma} f_{i\sigma}^\dagger c_{j\sigma} \quad (1)$$

and the local part, by

$$H_1 = U_{ff} \sum_i n_{if\uparrow} n_{if\downarrow} + U_{fc} \sum_{i, \sigma, \sigma'} n_{if\sigma} n_{ic\sigma'} + \epsilon_f \sum_i n_{fi} \quad (2)$$

I take the c -band centered around $E = 0$ and consider $U_{ff} = \infty$ (valid for f shells), so the f electrons are projected fermions, $X_{if\sigma} = (1 - n_{if-\sigma})f_{i\sigma}$, satisfying $[X_{i\sigma}, X_{j\sigma'}^\dagger]_+ = \delta_{ij}\delta_{\sigma\sigma'}(1 - n_{if,-\sigma})$. Using the Gutzwiller approximation, $X_{if\sigma} = q_\sigma f_{i\sigma}$ with $q_\sigma = (1 - n_f)/(1 - n_{f\sigma})$ and $n_{f\sigma} = (1/N) \sum_i \langle f_{i\sigma}^\dagger f_{i\sigma} \rangle$. This implies $(t_f, \epsilon_f) \rightarrow q_\sigma^2(t_f, \epsilon_f)$ and $V_{fc} \rightarrow q_\sigma V_{fc}$ in what follows, and corresponds to the slave-boson mean-field theory (SB-MFT), yielding a narrow, coherent f band with

a width $W \simeq k_B T_K^{mf}$ [6], the mean-field Kondo scale. Consistent with LDA calculations [14], I take $\epsilon_f \simeq E_F$. Below, I investigate the fate of this SB-MFT Kondo scale in presence of strong, quantum fluctuations of the f occupation, caused by the competition between mean-field coherence, $T_K^{mf}(V_{fc}, (t_{f,p}/U_{ff}))$ and incoherence, driven by U_{fc} .

I start by splitting the V_{fc} term as $(V_{fc} - \sqrt{t_f t_p}) \sum_{\langle i,j \rangle, \sigma} (f_{i\sigma}^\dagger c_{j\sigma} + h.c.) + \sqrt{t_f t_p} \sum_{\langle i,j \rangle, \sigma} (f_{i\sigma}^\dagger c_{j\sigma} + h.c.)$ and consider $H = H_0 + H_1$ with $V_{fc}^{(1)} = \sqrt{t_f t_p}$ to begin with. Using $a_{i\sigma} = (u f_{i\sigma} + v c_{i\sigma})$, $b_{i\sigma} = (v f_{i\sigma} - u c_{i\sigma})$ with $u = \sqrt{t_f/(t_f + t_p)}$, $v = \sqrt{t_p/(t_f + t_p)}$. it is easy to see that $H = H_0 + H_1$ is $H_0 = -t \sum_{\langle i,j \rangle, \sigma} (a_{i\sigma}^\dagger a_{j\sigma} + h.c.)$ and $H_1 = U_{fc} \sum_{i,\sigma,\sigma'} n_{ia\sigma} n_{ib\sigma'} + \epsilon_f \sum_{i,\sigma} [n_{ia\sigma} + n_{ib\sigma} + (a_{i\sigma}^\dagger b_{i\sigma} + h.c.)]$.

This is the spin $S = 1/2$ Falicov-Kimball model (FKM) with a *local* hybridisation term, which is finite whenever $\epsilon_f \neq 0$. Remarkably, when $\epsilon_f = 0$, this reduces to the pure $S = 1/2$ FKM! Below, we show how ($\epsilon_f = 0, V_{fc} = \sqrt{t_f t_p}$) separates *two*, different metallic phases.

First, at $\epsilon_f = 0$, we see that $[n_{ib}, H] = 0$ for *each* i , implying a *local* $U(1)$ invariance of H : local configurations with $n_b = 0, 1$ are rigorously degenerate. This is *exactly* the condition for having singular b -“number” fluctuations. As is known [15], the symmetry unbroken metallic phase is consequently *not* a FL, but is dominated by a superposition of one-particle ($n_b = 0$) and two-particle ($n_b = 1$) states at low energy. The a -fermion propagator is, assuming a lorentzian unperturbed DOS with half-width W for analytical clarity, very simple, showing “upper” and “lower” Hubbard bands, with a pseudogap at $E_F(=0)$:

$$\rho_a(\omega) = \frac{1 - n_b}{\omega^2 + W^2} + \frac{n_b}{(\omega - U_{fc})^2 + W^2} \quad (3)$$

The corresponding b -fermion propagator has branch cut singular behavior precisely at $E_F(=0)$, leading to singularities in the local, one- and two-particle responses: $\rho_b(\omega) \simeq \theta(\omega)|\omega|^{-(1-\alpha_0)}$ and $\chi_{ab}''(\omega) = \int dt e^{i\omega t} \langle a_{i\sigma}^\dagger b_{i\sigma}(t); b_{i\sigma}^\dagger a_{i\sigma}(0) \rangle \simeq \theta(\omega)|\omega|^{-(2\alpha_0 - \alpha_0^2)}$. Here, $\alpha = (1/\pi) \tan^{-1}(U_{fc}/W)$ is the so-called s -wave phase shift of the Anderson-Nozieres-de Dominicis (AND) X-Ray Edge (XRE) problem. Obviously, the FL quasiparticle weight, $Z = 0$, and the ω, T dependence of physical quantities will be governed by power-law responses. Notice that, in the (f, c) basis, divergence of $\chi_{ab}''(\omega)$ corresponds to extended (cf. non-local hybridisation), singular quantum fluctuations of the f occupation: it is precisely these fluctuations which destroy FLT at $\epsilon_f = 0$ via the AND orthogonality catastrophe (OC) [16]. Given that the $4f_{7/2}$ level hybridizes with *two* “ c ” bands in reality [14], we get the OC exponent, $\alpha = 2\alpha_0$.

In our FKM with $\epsilon_f = 0$, and at $T = 0$, the total fermion number, $n = n_a + n_b = n_f + n_c$, jumps dis-

continuously from $n_- = (1/2) + (1/\pi) \tan^{-1}(U_{fc}/2W)$ to $n_+ = (3/2) - (1/\pi) \tan^{-1}(U_{fc}/2W)$ for a range of densities, n , near unity. This corresponds to a sudden jump in the b -occupation, n_b , giving a *first* order “valence” transition as ϵ_f is tuned through $\mu(=0)$. At finite T , this line of first-order transitions ends at a second order critical end-point (CEP), and n_b varies very rapidly over an energy scale $O(k_B T)$ around $\epsilon_f = 0$, extrapolating to a jump $T = 0$. Notice that this jump in n_b depends on U_{fc}/t , which we choose henceforth to be such that this jump is vanishingly small [17] at $T = 0$, giving a quantum critical end-point (QCEP). This explicitly shows the link between emergence of singular non-FL behavior and the selective Mott localization of the b -electrons. *We emphasize that, with finite t_f, V_{fc} , both, the f - and c -fermions remain mobile: only their combination, $b_\sigma = (v f_\sigma - u c_\sigma)$, is localized.* Recent slave boson approaches have proposed this “selective Mott transition” of f -electrons in the context of the QCPs in RE systems [6]. Our work is a concrete, DMFT-based, realization of the “selective Mott” QPT, with a *non-local* hybridization. In contrast to earlier work [6], nFL behavior here arises from the AND-OC in the corresponding impurity problem as V_{fc} is varied across a critical value, $V_{fc}^{(1)} = \sqrt{t_f t_p}$.

We now show how the unique observations at the FQCP in $YbRh_2Si_2$ are understood as a consequence of the AND-OC derived above. The singularity in the b -DOS implies that their contribution to thermodynamic responses dominates that of the “itinerant” a -fermions. Hence, the low- T specific heat is

$$C_{el}(T) \simeq T \lim_{\eta \rightarrow 0} \text{Im}[G_{bb}(\omega + i\eta)]|_{\omega=T} \simeq T^\alpha. \quad (4)$$

giving the γ co-efficient as $\gamma(T) \simeq T^{-(1-\alpha)}$; actually, goes like $T^{-(1-\alpha)} \log(T/T_{coh})$ as the QCP is approached. The log- factor comes from seeing that the DOS is approximately a lorentzian with a maximum varying like $T^{-(1-\alpha)}$ and a full width at half-maximum equal to $\alpha\pi T$. When $E_F(=0)$ lies within this peak, we can write $E_F(T) = -\alpha\pi T$, whence the asymptotic form of $C_{el} \simeq T^{-(1-\alpha)} \log(T/T_{coh})$ follows. The entropy is then directly obtained as $S(T) = \int_0^T \gamma(T') dT' \simeq T^\alpha$. With $\alpha_0 = 0.3$, we thus find that both $C(T), S(T)$ vary as T^α with $\alpha = 0.6$, in nice agreement with observations in the nFL regime as a function of T [2].

What about transport?. In the impurity limit (note that the lorentzian unperturbed DOS in DMFT will not modify the “impurity” result), we have $G_{a0}(\tau) = (\pi T \rho_0 / \sin(\pi T \tau))$ and $G_{b0}(\tau) = \text{sgn}(\tau)/2$, whence the respective self-energies are $\Sigma_a(\tau) = U_{fc}^2 G_{a0}(\tau) G_{a0}(\tau) G_{a0}(\tau)$ and $\Sigma_b(\tau) = U_{fc}^2 G_{a0}(\tau) G_{a0}(\tau) G_{b0}(\tau)$. Direct evaluation followed by Fourier transformation then gives $\Sigma_b(i\omega_n) = -i(U_{fc} \rho_0)^2 [\omega_n (\ln(E_F/T) - \Psi(\omega_n/2\pi T) - \pi T)]$ and $\Sigma_a(i\omega_n) \simeq (\omega^2 + \pi^2 T^2)$. The dc resistivity within

DMFT is then $\rho_{dc}(T) \simeq (m/ne^2)\text{Im}\Sigma_b(\omega)|_{\omega=T} \simeq AT$, i.e., it is linear in T . The optical scattering rate will also be linear in ω . This is exactly in accord with observations near the QCP in $YbRh_2Si_2$ [8, 9]. Interestingly, with log-singularities in Σ_b above (which already imply $Z = 0$), higher-order terms, which must be carefully examined, lead precisely to the *branch cut* singular structure [19] for $\rho_b(\omega)$, characteristic of the “lattice X-ray edge” problem found in DMFT.

To proceed, observe that the impurity model corresponding to H_{FKM} can be bosonized in each radial direction centered around the “impurity” site [20]. For general band-filling, $n = n_a + n_b \neq 1$ per site, the umklapp terms from U_{fc} are irrelevant and hence ignored. The bosonized Lagrangian then describes a collection of *non-interacting* charge- and spin density collective modes:

$$L'_0 = \sum_{\rho,\sigma} \frac{u_{\rho,\sigma}}{2} \int [K_{\rho,\sigma} \Pi_{\rho,\sigma}^2(r) + \frac{1}{K_{\rho,\sigma}} (\partial_r \phi_{\rho,\sigma}(r))^2] dr \quad (5)$$

and $L_0'' = \frac{g}{\pi u_\rho} \sum_\rho \int \partial_r \phi_\rho(r) dr$. Here, $g, u_{\rho,\sigma}, K_{\rho,\sigma}$ are explicit functions of U_{fc}/t . Thus, interactions simply “shift” the *charge* bosonic modes relative to their free values. Introducing the usual symmetric-antisymmetric (charge-spin) combinations of $\phi_{\rho,\sigma}(r)$, we see that the antisymmetric (spin) channel completely decouples from the charge channel: a kind of high-dimensional **spin-charge separation**! This has been strongly emphasized by Anderson [16] in the cuprate context, and has important consequences, detailed below.

$\epsilon_f \neq 0$ has two effects: (i) it moves the b -fermion level away from E_F , and, (ii) finite a - b hybridisation generates a finite, but heavy b -fermion mass, due to recoil in the XRE problem [21], giving a small “coherence scale”, $\epsilon_{rec} = k_B T_{coh}$, below which HFL behavior obtains in the lattice model. The FL quasiparticle overlap, $Z \simeq e^{-C(t=\infty)}$, with $C(t) = 2U_{fc}^2 \int \frac{\chi_{ab}(\omega)}{\omega^2} (1 - \cos(\omega t)) d\omega$. This gives $Z \simeq \exp[U_{fc}^2 (\ln(\kappa)/(1 - \kappa^2))]$. Here, $\kappa = m_a/m_b$, with m_a the band mass of the a -fermion and m_b the heavy mass of the b fermion. Hence $T_{coh} \propto Z$ (note the difference from the SBMFT scale, T_K^{mf}) increases with ϵ_f , as indeed observed in the region to the right of the FQCP. In $D = \infty$, the relevant hybridization, $\epsilon_f \sum_{i,\sigma} (a_{i\sigma}^\dagger b_{i\sigma} + h.c)$ implies that the one-electron DOS will show a narrow, low-energy FL resonance, with upper/lower Hubbard bands at high energies, as is known [15]. Away from $V_{fc} = \sqrt{t_f t_p}$, the term $\sum_{\langle i,j \rangle, \sigma} \delta V_{fc} (f_{i\sigma}^\dagger c_{j\sigma} + h.c)$ also causes one-particle *intersite* hybridisation between the a, b fermions. This again gives the b -fermions a finite mass and results in another HFL with $T_{coh} \propto Z \ll 1$. Low-energy responses are then those of a HFL with $Z \ll 1$: an enhanced $\gamma = C_{el}(T)/T$, $\chi(T) = \chi_0$, $\rho_{dc}(T) = \rho_0 + AT^2$, etc, followed by a *smooth* crossover to the non-FL response found for $V_{fc}^{(1)} = \sqrt{t_f t_p}$, $\epsilon_f = 0$.

Using the bosonized form, Eq.(7), above allows further progress in the nFL regime. Expressing the transverse spin correlation function as an average over the phase variables permits its evaluation using $L_{0,\sigma}$. The result, following [22], is

$$\chi^{+-}(\mathbf{q}, \omega) = \frac{A}{TK_\sigma^{-1} - K_\rho} F\left(\frac{\omega}{T}\right) \quad (6)$$

with $K_\rho = \sqrt{v_F/(v_F + U_{fc})}$ and $v_F = 2t$ the Fermi velocity. $F(x)$ is a scaling function, $\simeq x$ for $x \ll 1$ and $\simeq 1$ for $x \gg 1$. And $K_\sigma = 1$ for the SU(2) invariant case, but $K_\sigma < 1$ including spin-orbit ($s - o$) coupling effects. This immediately yields the power-law T -dependence of the NMR relaxation rate as

$$\frac{1}{T_1} = \frac{T}{\omega} \sum_{\mathbf{q}} \text{Im} \chi^{+-}(\mathbf{q}, \omega) \simeq T^{-(K_\sigma^{-1} - K_\rho)} \quad (7)$$

The uniform spin susceptibility follows as $\chi(\mathbf{q} = 0, T) \simeq T^{-(K_\sigma^{-1} - K_\rho)}$. In DMFT, the singular-in- ω part of $\chi(\mathbf{q}, \omega)$ is independent of \mathbf{q} . This explains why *both* $T_1^{-1}(T)$ and Knight shift, $K_s(T)$, scale like $T^{-(K_\sigma^{-1} - K_\rho)}$ like $\chi(\mathbf{q} = 0, T)$. With a reasonable choice of U_{fc}/t , we get $K_\rho = 0.4$, leading to very good agreement with the host of power-law behaviors found in the magnetic response near the FQCP in $YbRh_2Si_2$. In particular, with the choice $K_\sigma = 1$, $\chi(\mathbf{q} = 0, T), T_1^{-1}(T), K_s(T)$ all follow a $T^{-0.6}$ law. Further, taking $\alpha_0 = 0.3$ (see above), we find $\chi/\gamma \simeq T^{-0.2}$. Assuming b/T scaling, where $b = (B - B_c)$ is the distance from the critical field, this implies $\chi(b)/\gamma(b) \simeq b^{-0.2}$, and that $\chi(b) \simeq b^{-0.6}, \gamma(b) \simeq b^{-0.4}$ near the FQCP. Further, with $A(b) \simeq 1/b$ [2], we find that the Woods-Saxon ratio, $A/\gamma^2 \simeq b^{-0.2}$ and $A/\chi^2 \simeq b^{0.2}$, which is weakly b -dependent and saturates at “higher” b . All these are experimentally seen [12]. While $\alpha_0 = 0.3$ as found in our *model* DMFT may change somewhat in a truly “first principles” theory, the qualitative theory-experiment agreement is compelling.

The higher- D spin-charge separation implied by the bosonized form of the impurity model also leads to consistency with the magnetotransport results: the Hall relaxation rate is now controlled by spinon-spinon scattering, leading to $\cot\theta_H(T) \simeq c_1 T^2 + c_2$ [16]. With $\rho_{dc}(T) = A_c T$, we get the Hall resistivity, $\rho_{xy}(T) \simeq T^{-1}$. Both are in good agreement with experiment.

We now turn to the evolution of the FS across the FQCP in $YbRh_2Si_2$. Hall data suggest an abrupt reconstruction of the FS across the FQCP as $T \rightarrow 0$. Thus, a large FS in the HFL regime, also seen in dHvA work [23], abruptly goes over to a small FS on the AF side. Within DMFT, in the symmetry-unbroken metallic phase(s), the shape *and* size of the FS is *not* affected by interactions, since the self-energy is purely local: $\Sigma_{a,b}(k, \omega) = \Sigma_{a,b}(\omega)$. Exactly at the FQCP, i.e., at $V_{fc} = \sqrt{t_f t_p}$, $\epsilon_f = 0$, the FS

will be a single sheet with a volume corresponding to the a -fermion number. To the right of the FQCP, the finite $\delta V_{fc}, \epsilon_f$ gives a finite $a - b$ hybridization, giving a HFL metal, as found above. In this regime, DMFT studies on the PAM with a relevant hybridization yield a large FS containing *both*, the lighter a - as well as the heavy b -fermions. Thus, the abrupt change in the FS volume is intimately linked with the selective Mott localization of the $b_\sigma = (vf_\sigma - uc_\sigma)$ fermions; they “decouple” from the FS at the FQCP. This occurs exactly at the point where the FL coherence scale vanishes, giving a non-FL metal with low-energy singular responses.

What about AF order? For small $\delta V_{fc}, \epsilon_f \ll 1$, *two-particle* processes, generated to second order in $\delta V_{fc}, \epsilon_f$, are more relevant than the one-particle $a - b$ hybridization. These processes couple two “impurities”, and lead to two-particle instabilities, as in coupled $D = 1$ Luttinger liquids [22]. To this order, extra terms, $H_{res}^{(2)} \simeq -\lambda^2 \sum_{\langle i,j \rangle} a_{i\sigma}^\dagger b_{i\sigma} b_{j\sigma'}^\dagger a_{j\sigma'}$, with $\lambda^2 \simeq O((\delta V_{fc})^2/U_{fc})$, are generated in H . In $D = \infty$, these are decoupled as $H_{res}^{eff} = -\lambda^2 \sum_{\langle i,j \rangle, \sigma} (M_{ab} a_{i\sigma}^\dagger b_{i\sigma} + M_b a_{i\sigma}^\dagger a_{j,-\sigma} + h.c.)$. Solving $H = H_{FKM} + H_{res}^{(2)}$ within DMFT should yield an AF metallic phase [18]: it will have the same symmetry as the AF phase in an “itinerant” view, since *both* have $M_a = \langle a_{i\sigma}^\dagger a_{j-\sigma} \rangle > 0$. However, I choose a different route. Bosonizing these terms in H_{res}^{eff} , the second term, corresponding to AF order, generates a cosine term in the bosonized Lagrangian for the spin sector: $L_\sigma^{int} = g_1 \cos(\beta \phi_\sigma)$, with $\beta = \sqrt{8\pi K_\sigma}$. With spin-orbit interaction, $K_\sigma < 1$, and the Lagrangian in the spin sector,

$$L_\sigma = L_{0,\sigma} + g_1 \int \cos(\beta \phi_\sigma) dr \quad (8)$$

is a quantum sine-Gordon model with a relevant cosine term. This leads to an AF ordered state, corresponding to a finite expectation value of the ϕ_σ field: $\langle \phi_\sigma \rangle > 0$ [22]. Thus, AF order here results as a particle-hole instability of the singular, non-FL metal derived above in the spin channel, rather than from a band FS instability, as would be the case in “conventional” cases where FS nesting features in a FL metal give itinerant magnetism: the latter picture cannot account for power-law responses seen at the QCP in $YbRh_2Si_2$. Thus, in our EPAM, at the FQCP, selective b -fermion localization permits AF to arise simply due to “inter-impurity” (corresponding to onset of RKKY-like) $b - b$ local moment correlations induced via “itinerant” a -fermions. Obviously, the FS now has a small volume, containing only the “itinerant” a -fermions.

At $V_{fc}^{(1)}$, the AND-OC will *always* occur in the symmetry-unbroken metallic phase in *any* $D < \infty$. Thus, we expect that our findings will survive inclusion of non-local correlations beyond DMFT. Also, the f -electrons

are *never* strictly localized: only the b -combination localizes at $V_{fc}^{(1)}$. An ab-initio theory for $\alpha = 0.6, K_\rho = 0.4$ and K_σ used here is hard: here, we have employed plausible $U_{fc}/t = 10$ (this is the *only* free parameter in our model) values. A truly first-principles correlated program (e.g. LDA+DMFT) is required to *derive* them. We plan to address this issue in future.

In conclusion, a *local* QCP, triggered by the AND-OC [24], is found in the DMFT solution of the EPAM as the model parameters are varied. Using high- D bosonization, non-FL responses with an uncanny resemblance to those found at the FQCP in $YbRh_2Si_2$ are uncovered. This QCP is unstable, either to a HFL, or to AF. *All* these findings are in very good qualitative agreement with the $T - b$ phase diagram of $YbRh_2Si_2$, whose unconventional QCP is thence proposed to be of the local type, and associated with the selective Mott localization in the EPAM. Our analysis is potentially applicable to other, d - and f -electron based systems showing non-FL behaviors near the $T \rightarrow 0$ itinerant-localized transitions.

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- [23] P. M. C. Rourke *et al.*, cond-mat/0807.3726, where a “large” high-field FS is claimed from detailed dHvA studies. By continuity, it should remain large in the entire HFL regime.
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